arXiv:0912.3320v2 [hep-ph] 3 Apr 2010

The INTEGRAL/SPI 511 keV Signal from Hidden Valleys in Type Ia and Core Collapse Supernova Explosions

Hooman Davoudiasl*[1] and Gilad Perez†[2]

[1]*Department of Physics, Brookhaven National Laboratory, Upton, NY 11973, USA*

[2]*Department of Particle Physics & Astrophysics, Weizmann Institute of Science, Rehovot 76100, Israel*

Abstract

We examine under what circumstances the INTEGRAL/SPI 511 keV signal can originate from decays of MeV-scale composite states produced by: (A) thermonuclear (type Ia) or (B) core collapse supernovae (SNe). The requisite dynamical properties that would account for the observed data are quite distinct, for cases (A) and (B). We determine these requirements in simple hidden valley models, where the escape fraction problem is naturally addressed, due to the long lifetime of the new composite states. A novel feature of scenario (A) is that the dynamics of type Ia SNe, standard candles for cosmological measurements, might be affected by our mechanism. In case (A), the mass of the state mediating between the hidden sector and the SM e^+e^- could be a few hundred GeV and within the reach of a 500 GeV e^+e^- linear collider. We also note that kinetic mixing of the photon with a light vector state may provide an interesting alternate mediation mechanism in this case. Scenarios based on case (B) are challenged by the need for a mechanism to transport some of the produced positrons toward the Galactic bulge, due to the inferred distribution of core collapse sources. The mass of the mediator in case (B) is typically hundreds of TeV, leading to long-lived particles that could, under certain circumstances, include a viable dark matter candidate. The appearance of long-lived particles in typical models leads to cosmological constraints and we address how a consistent cosmic history may be achieved.

* email: hooman@bnl.gov

† email: Gilad.Perez@weizmann.ac.il

I. INTRODUCTION

The recent INTEGRAL/SPI measurement of the 511 keV signal [1] from e^+e^- annihilation in the Milky Way reconfirms [2] its puzzling characteristics: it appears that e^+e^- pairs annihilate at a rate of $1.5 \pm 0.1 \times 10^{43}\,\mathrm{s}^{-1}$ in the Galactic bulge, whereas the corresponding rate in the disk is $0.3{\pm}0.2{\times}10^{43}\,\mathrm{s}^{-1}$. Hence, the bulge-to-disk annihilation ratio $B/D \sim 3{-}9$ is implied [1]. The size of the signal and the inferred value of B/D make it a challenge to explain the data using standard astrophysical sources.

A leading conventional candidate would be the radioactive β^+ decays of type Ia Supernova (SN) products, such as ^{56}Co. These explosions can supply the necessary flux of e^+ and are thought to happen mostly in the central region of the galaxy, with an old stellar population [1]. However, it has been estimated that the fraction of the e^+ that can escape the type Ia ejecta is not large enough to produce the entire observed signal [3].

It has also been argued that the bulge rate of SN type Ia ($\mathrm{SN_{Ia}}$) explosions can be inferred by mass-scaling of the rate in early type galaxies that host old stars [4]; this would then lead to values of B/D that are too small [1] (however, this extrapolation ignores the differences between the evolutions of the Milky Way bulge and such galaxies). In this case, unless one assumes that e^+e^- annihilation takes place far (on galactic scales) from the source [4], this ratio cannot be straightforwardly explained.

Various other sources, such as low mass X-ray binaries [1], have been proposed to explain the 511 keV signal. However, generally speaking, all such explanations suffer from rather sizable uncertainties and it is not clear if any plausible astrophysical mechanism can account for the data. This situation has motivated new physics proposals, such as those based on the Galactic dark matter (DM) content; see for example [5–7].

Here, we study a different set of new physics possibilities that do not depend on cosmic relic abundances of particle populations, but rely only on supernova production of new composite states that decay into e^+e^-. We consider two distinct scenarios, referred to hereafter as cases (A) and (B):

- In case (A), the new states are produced in type Ia explosions.
- In case (B) we consider production in core collapse (cc) supernovae.

A simple estimate shows why this may, in principle, be energetically possible. Roughly

speaking, at a rate of 1–2 explosions per century, $\mathcal{O}\left(10^{48-50}\right)$ MeV s^{-1} of binding energy is released into the Galaxy. Obviously, even if a tiny fraction of this energy flux is carried by the hidden sector, it would be enough to account for the e^+e^- annihilation signal. In fact, it was proposed long ago that the 511 keV signal could have been generated by the decays of heavy τ-neutrinos (now ruled out) produced in cc SN ($\mathrm{SN_{cc}}$) explosions [8].

As mentioned above, the 511 keV data could have reasonable standard explanations [9]. Here, we do not attempt to refute possible standard astrophysical descriptions of the data. However, given the uncertainties involved in these (or other) explanations, we believe alternative proposals merit attention. In particular, the long lifetime of particles in the scenarios we consider can naturally address the escape fraction problem present in the conventional $\mathrm{SN_{Ia}}$ picture. Also, our models could have testable predictions for collider experiments [case (A)], and may also provide DM candidates [case(B)], though we do not *require* that our scenarios give rise to a viable DM particle. We note in advance that, while the new scenarios we examine provide advantages over the conventional accounts, they are subject to their own challenges and constraints. Nonetheless, they offer concrete and interesting examples of how hidden valley dynamics could affect astrophysical phenomena, such as the Galactic 511 keV signal.

The 511 keV signal suggests that $\sim 10^{43}$ e^+e^- pairs are annihilated each second, within the bulge at the center of the Galaxy. Astronomical observations require that e^+ be injected into the interstellar medium (ISM) near threshold, at energies below a few MeV per e^+ [10]. We will propose that this injected flux comes from the decay of a new state X of mass $m_X \sim 1$ MeV, from a hidden dynamical sector of the type discussed in Ref. [11]. We further assume that X only decays through SM final states and its branching ratio into e^+e^- is typically $\mathcal{O}(1)$. Given the complexity of the astrophysical system and the fact that the hidden dynamics is non-perturbative our analysis is only done in a semi-quantitative manner. Yet, the analysis presented below is concrete enough to make the point that hidden valley models may play an important role in dynamics of astrophysical systems and the cosmological evolution of the universe.

In case (A), we find that a successful scenario based on $\mathrm{SN_{Ia}}$ requires a mediator of mass $\mathcal{O}(300)$ GeV, connecting the SM and the hidden sectors, making its discovery at a 500 GeV e^+e^- collider a possibility. We also show how present experimental constraints may be addressed. Here, we also briefly discuss how kinetic mixing of the photon with a light vector

can provide another potential mediation mechanism between the hidden and the SM sectors [12]. The production of the X particle is Boltzmann suppressed, hence it is not relativistic and its e^+e^- decay products have MeV-scale energies, as required by observations. The lifetime of X is long enough to exit the SN ejecta before decaying and avoids the conventional $\mathrm{SN_{Ia}}$ positron escape fraction problem. The $\mathrm{SN_{Ia}}$ Galactic population is centrally peaked and helps explain the spatial distribution of the 511 keV signal. Potentially troublesome long-lived particles produced in the early universe can be efficiently annihilated via strong dynamical processes within the hidden sector. A major concern in this scenario, apart from direct constraints on the model, is whether Big Bang Nucleosynthesis (BBN) can go through without much perturbation. Although we show that this may be plausibly achieved within the considered framework, perhaps a more detailed analysis is required to reach a firm conclusion.

In case (B), the X production in $\mathrm{SN_{cc}}$ is not Boltzmann suppressed and can be brought down to the required levels by lowering the coupling of the hidden and SM sectors. This is achieved by assuming a high SM-hidden mediation scale of $\mathcal{O}(100)$ TeV, which leads to a long lifetime for X. This feature allows the e^+e^- decay products to be deposited far from the $\mathrm{SN_{cc}}$ sources that are mostly found in the Galactic disk. This long decay length could lead to acceptable B/D values, since it causes the surface brightness of the disk e^+e^- annihilations to be reduced and may also help with magnetic transport of some positrons into the bulge. Even though the kinetic energy of the initially produced hidden "partons" is $\mathcal{O}(10)$ MeV, fragmentation into X and other hidden hadrons softens the final state e^+e^- energy to the desired few MeV level. Here, acceptable cosmic abundance of long-lived particles may require low reheat temperatures, but simple models give rise a to a plausible DM candidate. The main shortcoming in this scenario stems from the fact that the disk population of the $\mathrm{SN_{cc}}$ is bigger than that of the bulge and a Galactic mechanism is needed to transport the generated positrons into the bulge, as we will discuss.

In the following section, we discuss the microscopic structure of our models. In section III, we explain how the 511 keV signal can be addressed via type Ia supernova explosions. Section IV deals with the core collapse case. In section V, we discuss the cosmological implications of our setup. Section VI contains our final conclusions.

II. MICROSCOPIC SETUP

We assume that the hidden sector has a dynamical scale $\Lambda_h \sim 1$ MeV at which the new states emerge. This scale sets the masses of the new hidden hadrons[1]. The decays of the composite states then yield the requisite e^+ input at $\mathcal{O}$(MeV) energies . Of course, if X is emitted at energies much larger than its mass, the positrons will be very relativistic, in conflict with observational evidence that demands a cold positron flux. However, below we will argue that Boltzmann suppression, in case (A), or parton shower and fragmentation, in case (B), makes this unlikely, given the relevant energies. We will denote the fermionic "quark" degrees of freedom in the hidden sector, subject to non-trivial hidden dynamics, by f. Below the scale Λ_h, strong dynamics leads to the emergence of hidden hadrons made up of f quarks.

Following Ref. [11], we will assume that the hidden dynamics is coupled to the SM via higher dimension operators suppressed by a scale M. For concreteness, we will consider the following dimension-6 operator that can arise in a variety of models:

$$\frac{(\bar{\psi}_i \Gamma_1 \psi_j)(\bar{f}_a \Gamma_2 f_b)}{M^2}, \tag{1}$$

where ψ_i is an SM fermion and $\Gamma_{1,2}$ represent a product of Dirac matrices. For example, if the visible and the hidden sectors couple through a massive scalar or vector Φ of mass m_Φ, then one gets $M \sim m_\Phi/\sqrt{g_1 g_2}$ [11], where g_1 and g_2 denote couplings of the Φ to the SM and hidden sectors, respectively. In what follows, we focus on the case where Φ is a vector particle. In our discussion of case (A), we will briefly consider how the kinetic mixing of a light vector particle with the photon [12] can provide another mediation mechanism between the hidden and the SM sectors.

In places where a concrete model is appropriate to consider we will adopt the One-Light-Flavor (1LF) model discussed in Ref. [11]. We will denote the lightest pseudo-scalar by η_h and the lightest vector by ω_h, with $m_{\omega_h} \sim \sqrt{n_c^h}\, m_{\eta_h}$ [11], based on the results from Ref. [13]; n_c^h is the number of hidden colors. There is also a scalar, σ_h, that is typically intermediate in mass and lifetime [11]. The generic expectation is that σ_h is up to 30% heavier than η_h [11],

[1] For simplicity, we do not consider the possibility of having very light pseudo-Goldstone bosons (PGBs) well below Λ_h. Inclusion of PGB states would further enrich the structure of the theory and would add new scales to the system, hence it is worth future study.

and hence it would not decay into $\eta_h + e^+e^-$, within our typical parameter ranges. Then the key aspects of the phenomenology we are interested in can be captured by considering only η_h and ω_h, as we do for simplicity in the following. We also briefly discuss below the possible implications of the $\sigma_h \to \eta_h\gamma\gamma$ loop-mediated decay.

III. CASE (A): TYPE IA SUPERNOVA EXPLOSIONS

Let us begin with an estimate of the X-production in $\mathrm{SN_{Ia}}$. These events are attributed to the explosion of accreting white dwarfs through thermonuclear reactions that disrupt the star entirely, releasing $\sim 1 \times M_\odot \approx 2\times 10^{33}$ g of binding energy in the process. Observations of similar processes in other galaxies, including those of novas that are believed to share the same progenitor, point to a centrally peaked distribution of $\mathrm{SN_{Ia}}$ [1]. It is then plausible to assume that they could give rise to $B/D \gtrsim 1$, as long as one can provide a way for a large enough fraction of the generated e^+ flux to escape the explosion and reach the interstellar medium (ISM). We will base our estimates on the carbon deflagration model of $\mathrm{SN_{Ia}}$, a detailed discussion of which can be found in Ref. [14]; we will adopt the parameters of the widely used W7 model.

The $\mathrm{SN_{Ia}}$ explosion occurs over about 1 s, during which a large fraction of the white dwarf mass, roughly $0.7M_\odot$, reaches temperatures of order $T_{\rm Ia} \approx 6\times 10^9$ K. We will assume a mean density $\rho \approx 3\times 10^8\, \mathrm{g\dot{c}m^{-3}}$ [14]. For the electrons in the white dwarf, bremsstrahlung processes in which an electron scatters from the electromagnetic filed of the nucleus N are the dominant process: $eN \to eNX$ [15]. The rate for scalar or vector emission from non-relativistic and non-degenerate plasmas has been calculated in Ref. [16]. Since $T_{\rm Ia} \lesssim m_e$, and degeneracy is mild for our reference parameters [15], we adapt their result for a reasonable estimate of energy loss rate, in erg $\mathrm{g^{-1}}$ $\mathrm{s^{-1}}$, [15, 16]

$$\varepsilon_X \approx e^{-m_X/T_{\rm Ia}}\alpha'\,\eta\, 2.8\times 10^{26}\, T_8^{0.5}\, Y_e\, \rho \sum_j \frac{X_j Z_j^2}{A_j}, \tag{2}$$

where $\alpha' \equiv g'^2/(4\pi)$, with g' the coupling of the scalar (massive vector) X state to the electron, corresponding to $\eta = 1(3)$. Here, $T_8 \equiv T/(10^8\,\mathrm{K})$, $Y_e \simeq 0.5$ is the electron number fraction relative to baryons, X_j is the mass fraction of species j with atomic number Z_j and mass A_j; we will take $\sum_j X_j Z_j^2/A_j \sim 14$ for our estimates. In the above formula, we have introduced a factor for Boltzmann suppression, since $T_{\rm Ia} < m_X$. We require that the above

rate yield $\sim 10^{43}$ X particles per second, in order to generate the inferred e^+e^- annihilation signal. Taking the SN_{Ia} rate to be roughly 1 per century in the Galaxy and $m_X = 2.0\,\text{MeV}$, we get

$$\alpha' \sim 4.6 \times 10^{-22}/\eta\,. \tag{3}$$

With $\eta = 3$ for a vector X and a naive estimate $g' \sim m_X^2/M^2$, we find $M \sim 300\,\text{GeV}$; here and throughout this work we will assume that the decay constant of a hidden hadron $f_X \sim m_X$.

A rough estimate of the X lifetime in this scenario is

$$\tau_X \sim \frac{16\pi}{m_X g'^2} \sim \frac{16\pi M^4}{m_X^5} \sim 10\,\text{s}\left(\frac{M}{300\text{ GeV}}\right)^4\left(\frac{2\text{ MeV}}{m_X}\right)^5\,. \tag{4}$$

The size of τ_X also provides a way to address the problem of e^+ escape fraction. To see this, we argue that the X particles will typically fly out with a velocity v_X, much faster than the stellar explosion which moves at a relatively low speed $v_{\text{Ia}} \sim 0.03$ [14]. The reason for this is that the energy distribution of the emitted X particles inherits the electron energy distributions inside the SN_{Ia} , hence we expect

$$v_X \sim (T_{\text{Ia}}/m_X)^{1/2} \gg v_{\text{Ia}}\,. \tag{5}$$

Finally, we note that the size of the progenitor of a SN_{Ia} is given by the radius of the accreting white dwarf, $R_{\text{wd}} \sim 10^4$ km ~ 0.03 s. Then, Eq. (4) implies that the X particles will decay well outside the stellar explosion front and will not be absorbed by the ejecta. This avoids the escape fraction problem of the conventional SN_{Ia} picture, where the positrons can get absorbed efficiently as they are released within the outward moving stellar matter.

A. Direct constraints

The mass scale M may seem unacceptably low, in light of the existing LEP data [17–19], or direct detection data [20]. However, we argue that simple assumptions about the nature of the coupling of the hidden and SM sectors to the mediating particle Φ (such as a heavy Z') can allow for such values of M. In particular, if we assume that the hidden sector couples to Φ strongly with $g_2 \sim 10$, whereas the SM couples to Φ weakly with $g_1 \sim 1/g_2$, then many existing bounds can be satisfied for $m_\Phi \sim 300$ GeV. As we do not require X production from SM hadrons in SN_{Ia} , a simple assumption would be to take the Z' coupling

to quarks to be negligible, in which case the Tevatron bounds [20] would not apply. Here we would like to mention that kinetic mixing of the photon with a light vector particle can offer another interesting possibility. For example, if the mixing is governed by a loop-level mixing parameter $\varepsilon \sim 10^{-3}$, for $g_2 \sim 0.1$ and $m_\Phi \sim 1$ GeV, we may achieve the required level of coupling between the hidden and the SM sectors. Such parameters can accommodate precision measurements of $(g-2)_e$ [21], while potentially accounting for a possible deviation of $(g-2)_\mu$ from the SM prediction [22]. However, for concreteness we will concentrate on the case with $m_\Phi \sim 300$ GeV, in what follows.

The LEP bounds on $(e^+e^-)_{V+A} \to (e^+e^-)_{V+A}$, from Ref. [17], roughly require $\Lambda \gtrsim 7-9$ TeV (depending on the sign of the effective dimension-6 operator), where $\Lambda = \sqrt{2\pi}\, m_\Phi/g_1$. For $g_1 \approx 1/10$ and $m_\Phi \approx 300$ GeV we then get $\Lambda \approx 8$ TeV which is in the realistic regime. We hence adopt a minimal case where Φ couples dominantly to $(e^+e^-)_{V+A}$ in the SM. We note that this would introduce a suppression by a factor of 1/2 in Eq. (2), due to the random polarization of the electrons in the star. Even then, $M \sim 300$ GeV still yields about 7×10^{42} e^+e^- pairs per second which, within the accuracy of our analysis, is of the right size.

Another possible LEP bound comes from $e^+e^- \to \gamma + \not{E}$ [18, 19], where the missing energy is carried by the invisible $\bar{f}f$ jets. We estimate the cross section for $e^+e^- \to \gamma \bar{f} f$ by

$$\sigma_{ef\gamma} \approx \frac{1}{4} \times \frac{\alpha}{\pi} \times \frac{n_c^h E_{\rm cm}^2}{12\pi M^4}\,, \tag{6}$$

where the factor of (1/4) accounts for chiral couplings of the initial and final states, the factor of α/π accounts for the initial state radiation, and $E_{\rm cm}$ is the center of mass energy. In deriving the above formula, we have simply adapted similar results, say, for $e^+e^- \to \mu^+\mu^-$ [23]. An analogous process arises in models with large extra dimensions [24], where the missing energy is carried off by Kaluza-Klein gravitons [25, 26]. Using the results of Ref. [26], we then infer that $M \approx 300$ GeV is a safe suppression scale, for $E_{\rm cm} \approx 200$ GeV and $n_c^h = 2$, given the LEP2 bounds [18, 19] on $e^+e^- \to \gamma + \not{E}$.

B. Astrophysical constraints

The above inferred value of α' is not constrained by astrophysical processes at $T \ll m_X$, such as those of red giant cores at $T \sim 10\,$keV. However, $\mathrm{SN_{cc}}$ explosions, characterized by $T \sim 30\,$MeV, could overproduce the 511 keV signal and, in principle, lead to severe

constraints on our SN_{Ia}-based scenario. Remarkably, for $M \sim 300\,\mathrm{GeV}$, we roughly infer a mean free path $\lambda \sim M^4/T^5 \sim 10^2\,\mathrm{m}$ for the hidden sector f quarks (even though we are assuming that the hidden sector has negligible couplings to SM hadrons of the core, we may expect that such a hot environment provides a photon and e^+e^- thermal bath that can efficiently generate a population of f-quarks). Since the size of the hot SN core is $10-100$ km, we see that the f quarks are quite likely trapped and can only be surface-emitted. This avoids constraints coming from over-cooling of the SN_{cc} core, which would otherwise halt the explosion.

One may also worry that a trapped "hidden-sphere" will contain a substantial fraction of the explosion energy and overproduce the 511 keV signal. However, given the estimation of the X lifetime in Eq. (4) we claim that over-production is not a worry. The reason is that the progenitors of the SN_{cc} are red or blue giant stars of radius $\sim 10^2 R_\odot \sim 10^2$ s, much larger than the decay length of X. Hence, the e^+e^- decay products of the X particles will be absorbed well-within the SN_{cc} progenitor and will not contribute to an over-population of e^+ in the ISM. In fact, this effect might help to prevent the shock wave (which drives the cc explosion) from being stalled, a major outstanding problem in understanding SN_{cc}.

Regarding σ_h, we note that it is expected to be at most 30% heavier than η_h [11]. Hence, within our typical parameter range, it would decay into $\eta_h\gamma\gamma$, via a one-loop radiative process, with an approximate life time

$$\tau_\sigma \sim 2\times 10^{11}\,\mathrm{s}\left(\frac{0.4\,\mathrm{MeV}}{\Delta m_h}\right)^5\left(\frac{M}{300\,\mathrm{GeV}}\right)^4 , \tag{7}$$

where $\Delta m_h = m_{\sigma_h} - m_{\eta_h} \sim 0.4\,\mathrm{MeV}$. This is roughly consistent with the bound [15] produced by the Solar Maximum Mission (SMM) satellite [27] looking for a flash of photons emitted from the 1987a SN_{cc}, yet may require Δm_h to be a factor of few smaller.

IV. CASE (B): CORE COLLAPSE SUPERNOVAE

As noted before, SN_{cc} can also produce the requisite e^+ flux to account for the size of the 511 keV signal. However, as we will see, this can be done in a very different regime of hidden sector models, due to the much higher temperature of such explosions. The estimated rate of energy release from the SN_{cc}, for particles that free stream out of the core, can be obtained from that of $\nu\bar{\nu}$ emission from a non-degenerate stellar core, via $NN \to NN\nu\bar{\nu}$, with N a

nucleon [15]

$$\varepsilon_\nu = 2.4 \times 10^{17}\,\mathrm{erg\,g^{-1}\,s^{-1}}\,\rho_{15}\,T_{\mathrm{MeV}}^{5.5}\,. \tag{8}$$

Here, the density ρ_{15} is in units of $10^{15}\,\mathrm{g/cm^3}$ and the temperature T_{MeV} is in MeV. In the $\mathrm{SN_{cc}}$ core, $T \sim 30\,\mathrm{MeV}$ and the rate saturates at $\rho_{15} \approx 0.03$ [15], and we find $\varepsilon_\nu \approx 9.6 \times 10^{24}\,\mathrm{erg\,g^{-1}\,s^{-1}}$. The temperature of the core drops over a time scale of order a few seconds. Given the steep temperature dependence of the rate in Eq. (8), we estimate the total f-quark (hidden sector) output from the emission over the first few seconds, by rescaling the Fermi constant $G_F \simeq 1.2 \times 10^{-5}\,\mathrm{GeV^{-2}}$ of the SM weak interactions. For an object of mass $\sim 1.5 M_\odot \sim 3 \times 10^{33}\,\mathrm{g}$, the ratio R_f of f-quark output to the total energy budget $\sim 3 \times 10^{53}\,\mathrm{erg}$ is given by

$$R_f \sim 10^4 (M^2 G_F)^{-2}\,. \tag{9}$$

Based on our preceding discussion, $R_f \sim 10^{-6}$ is a reasonable estimate for the required output. We then get $M \sim 100\,\mathrm{TeV}$ as a natural scale for the operator in Eq. (1), if one wants to explain the 511 keV signal from $\mathrm{SN_{cc}}$ production.

Once the f-quarks are produced, they would "hadronize" on their way out of the SN core. Showering and fragmentation of the f-quarks will then distribute their emission energy among several X and other hidden hadrons. As a rough guide, we will consider the average multiplicity N_h of QCD hadrons in e^+e^- collisions. Given the typical SN temperature of about $30\,\mathrm{MeV}$ and $\Lambda_h \sim 1\,\mathrm{MeV}$, the data at $\sqrt{s} \sim 30\,\mathrm{GeV}$ will provide a good analogue, and we find $N_h \approx 20$ [28]. We then expect $\mathcal{O}(10)$ hidden hadrons to result from f-quark emissions at $T \sim 30\,\mathrm{MeV}$. As long as the initial emitted parton fragments into at least a few leading hadrons, their energies will typically be below $\sim 10\,\mathrm{MeV}$, which will eventually yield positrons at few MeV energies, as required by the data. This feature is a universal aspect of generic strong dynamics, regardless of the specific model. Looking at a compilation of measurements for center of mass energies between 14-44 GeV, by the Tasso collaboration [29], we find, indeed, that in $e^+e^- \to q\bar{q}$ processes less than 1% of the events contain a particle which carries 50% or more of the beam energy and the majority of final state hadrons are soft.

A reasonable fraction of these hidden sector hadrons must decay well-outside the $\mathrm{SN_{cc}}$ progenitor star, if one wishes to disperse the e^+e^- throughout the ISM. Since the typical radii of the collapsing stars is $\mathcal{O}(10^2 R_\odot)$, we require $\tau_X \gtrsim 100\,\mathrm{s}$. Note that the naive

expectation for the lifetime of X in this scenario is now

$$\tau_X \sim \frac{16\pi M^4}{m_X^5} \sim 10^{11}\ \mathrm{s} \left(\frac{M}{100\ \mathrm{TeV}}\right)^4 \left(\frac{2\ \mathrm{MeV}}{m_X}\right)^5 , \tag{10}$$

which is more than sufficient to satisfy this requirement. As we will discuss below, this long lifetime could be a useful feature in explaining the distribution of the 511 keV signal.

The inferred e^+e^- flux seems to come mainly from the Galactic bulge and is not evenly distributed throughout the Galactic disk. However, the $\mathrm{SN_{cc}}$ events are expected to be mostly concentrated in the disk. The Galaxy can be roughly described as a disk of radius $r_G \approx 15\,\mathrm{kpc}$. We approximate the central bulge of the Galaxy as a sphere with a radius $r_B \sim 3\,\mathrm{kpc}$. A reasonable lower bound on the fraction f_B of the cc events that come from the bulge region is obtained by the ratio of the disk area inside and outside the bulge [30]

$$f_B \gtrsim r_B^2/(r_G^2 - r_B^2) , \tag{11}$$

which yields $f_B \gtrsim 0.04$. One can also get estimates of the f_B using observational data on neutron stars and pulsars [31]. The distribution n_{cc} of $\mathrm{SN_{cc}}$ in our galaxy can be described, in galactocentric cylindrical coordinates (r, z, θ), by [31, 32]

$$n_{cc}(r) \propto r^\xi e^{-r/u} \left[0.79\, e^{-(z/0.212)^2} + 0.21\, e^{-(z/0.636)^2}\right] , \tag{12}$$

where ξ and u are fit parameters, and z is in kpc. The neutron star data yield $\xi = 4$ and $u = 1.25$ kpc, whereas the pulsar distribution implies $\xi = 2.35$ and $u = 1.53$ kpc [31]. We find that the first set of parameters yields $f_B \approx 0.04$, whereas the second set gives $f_B \approx 0.11$. In either case, we need to address the paucity of bulge sources.

Ref. [4] has suggested that if positrons can propagate away from the Galactic disk they will fill out a larger volume and reduce the surface brightness of the disk annihilation emission. Also, once the positrons get out to $z \gtrsim 3$ kpc, they will escape beyond the cosmic ray halo (CRH), where Galactic (poloidal) magnetic fields may transport them into the bulge. Ref. [4] argues that, if these conditions are satisfied, then B/D as low as ~ 0.5 can be consistent with the current INTEGRAL/SPI observations. However, a naive estimate of the escape fraction f_{esc} from the CRH, using slow-down and confinement time scales for positrons, yields $f_{\mathrm{esc}} \sim 0.1$ [4]. This is not sufficient to get the required level of disk e^+ population into the bulge, yet there may be room for enhancing f_{esc}, given the uncertainties involved. We will next discuss how in our scenario $f_{\mathrm{esc}} \gg 0.1$ can be naturally obtained. This would easily allow us to take advantage of the mechanism suggested in Ref. [4].

The long lifetime τ_X in Eq. (10) is helpful in releasing the requisite e^+e^- away from the cc explosion and the Galactic disk. Here, we give a rough estimate for the fraction of the e^+e^- trapped within the CRH. We only consider e^+e^- from decays of the X particles in our scenario. For simplicity, we approximate the CRH by a slab of thickness d and assume that the planar extent of the slab (Galactic disk) is large compared to other length scales. Furthermore, it is also assumed that the $\mathrm{SN_{cc}}$ explosions are confined to a thin layer (set by the thickness of the Galactic disk) in the middle of the slab, which is a reasonable approximation. The fraction of the X particles that decay within the slab is at most of order

$$F_X \sim \frac{2}{\pi}\int_0^{\pi/2} d\theta\,(1 - e^{-l/\tau_X}), \tag{13}$$

where $l = d/(2\cos\theta)$. Guided by Eq. (10), let us assume that $\tau_X \sim 10^{12}\,\mathrm{s}$, equivalent to a decay length of order $10\,\mathrm{kpc}$; this corresponds to $m_X \approx 2m_e$. For $d = 6\,\mathrm{kpc}$, we then find $F_X \sim 0.5$; for τ_X equivalent to $20\,\mathrm{kpc}$ (M slightly larger than $100\,\mathrm{TeV}$) we get $F_X \sim 0.3$. Hence, we see that for typical values of the lifetime τ_X, we can easily suppress the fraction of the X particles that decay within the CRH to $F_X \lesssim 0.5$. The long distance decays of X, compared to the Galactic disk scale, ensures that the e^+e^- is dispersed over a larger volume, which helps diminish the surface brightness of the disk e^+e^- annihilation emission [4]. In addition, this provides a natural mechanism for transporting the e^+e^- beyond the CRH, where they could be guided into the bulge by the large scale poloidal magnetic field. Hence, the success of case (B) in explaining the 511 keV signal depends on whether enough positrons can be transported into the bulge. Whether or not this condition is satisfied in our Galaxy is a question beyond the scope of our paper and poses a challenge to a scenario based on $\mathrm{SN_{cc}}$.

The above discussion outlines how hidden dynamics of the type discussed in Ref. [11] can in principle produce the gross features of the observed 511 keV signal. However, some questions, such as that of the consistency of this type of scenario with standard cosmology, can only be answered within more specific models. As a simple example, we will adopt the One-Light-Flavor (1LF) model discussed in Ref. [11] and mentioned above in section II. The key aspects of the phenomenology we are interested in can be captured by considering only η_h and ω_h and not σ_h, as we do for simplicity in the following.

V. COSMOLOGY

In the early universe, once the cosmic temperature T falls below the dynamical scale $\Lambda_h \sim 1\,\mathrm{MeV}$, various hidden hadronic states appear. The fast interactions amongst the hidden sector hadrons allow them to annihilate into the lightest state η_h, as the universe gets cooler [11]. If the lightest state has a sufficiently short lifetime, $\tau \ll 1\,\mathrm{s}$, then the hidden sector decays into the SM degrees of freedom in time for BBN. However, this is generally not the case in our type of scenarios. Also, a post-BBN hidden hadronic gas could be unacceptable, as such an ensemble reshifts like matter and would come to dominate the universe well-ahead of the standard epoch near $T \sim 3\,\mathrm{eV}$. Let us examine the status of the 1LF model regarding these questions. From the results in Ref. [11], the lifetimes of η_h and ω_h (assuming that the same type of expressions are valid for $m \sim 1\,\mathrm{MeV}$) have the following dependencies

$$\tau_{\eta_h} \propto M^4/\left(f_{\eta_h}^2 m_{\eta_h}^5\right) \quad \text{and} \quad \tau_{\omega_h} \propto M^4/m_{\omega_h}^5\,, \tag{14}$$

respectively; as mentioned before we will take $f_{\eta_h} \sim m_{\eta_h}$.

A. Case (A)

With $M \approx 300\,\mathrm{GeV}$, $n_c^h = 2$, and $m_{\eta_h,\omega_h} \sim 1.5, 2$ MeV, respectively, we find $\tau_{\eta_h} \sim 10^{21}$ s and $\tau_{\omega_h} \sim 3$ s. Here, ω_h has the kind of lifetime needed for X in case (A). However, η_h is not sufficiently long-lived to be a safe DM candidate [33] and if produced in equilibrium, could come to dominate over radiation during the BBN, which is not an acceptable outcome. To examine these questions, let us estimate the decoupling temperature of the hidden sector.

Above $T \sim m_X$, we can treat f quarks as free, and their production in the plasma is governed by their interactions with e^+e^-. The rate for this interaction is given by $\Gamma_{\mathrm{ef}} \sim n_e\,\sigma_{\mathrm{ef}}\,v$. Here, $n_e \sim T^3$, the cross section σ_{ef} is given by Eq. (6), but without the factor of α/π, the relative velocity $v \sim 1$, and $E_{\mathrm{cm}} \sim 2T$. The process $e^+e^- \to f\bar{f}$ decouples when $\Gamma_{\mathrm{ef}} < H$, with the Hubble rate $H \simeq 1.7\, g_*^{1/2}\, T^2/M_{\mathrm{Pl}}$; g_* is the number of relativistic degrees of freedom and $M_{\mathrm{Pl}} \simeq 1.2 \times 10^{19}\,\mathrm{GeV}$ is the Planck mass. The decoupling temperature, for $g_* \simeq 10.8$ ($T \gtrsim m_X$) is then given by $T_d \approx 4\,\mathrm{MeV}$. Here, T_d is in principle high enough for successful BBN. A similar calculation also suggests that the X particle, identified in our 1LF example as ω_h, would be out of equilibrium with the SM plasma for $T \lesssim m_X$. A

simple assumption would be that the reheat temperature was in the 1-4 MeV range, such that the f-quarks did not equilibrate with the SM[2]. However, arranging for such a low reheat temperature would require a rather non-standard cosmic evolution (for bounds on the scale of low temperature inflation see *e.g.* [36] and references therein). We will next argue that the strong interactions amongst the hidden sector states could efficiently suppress the relic density of η_h through the relatively fast decays of ω_h, in equilibrium.

The $\eta_h - \omega_h$ system stays in equilibrium through fast hadronic interactions governed by a cross section $\sigma v \sim 1/m_{\omega_h}^2$. As the ω_h population decays, its number density decreases as

$$n(T) \sim T^3 e^{-\Gamma/H}, \tag{15}$$

where $\Gamma \approx 2.1\times10^{-25}$ GeV is the width of ω_h, for $m_{\omega_h} \sim 2$ MeV. The fast thermal interactions decouple once $n\sigma v \sim H$, which yields $T_d \sim 0.1$ MeV, and hence

$$n(T_d)/T_d^3 \sim e^{-\Gamma/H(T_d)} \sim 10^{-20}. \tag{16}$$

This is a very rough estimate, but shows that the strong hidden dynamics can efficiently suppress the η_h number density to negligible levels. We note that a more detailed analysis may be called for to determine whether the proximity of the BBN era and the onset of hadronization and decays in the hidden sector does not cause large deviations from the standard Big Bang picture. It is important to note, however, that the X particle density just below the MeV temperature, when BBN starts, drops very rapidly with falling temperature, since $n(T) \propto e^{-\Gamma/H}$ and $\Gamma/H \propto m_X^5/T^2$. For instance, for $m_X = 3\,\text{MeV}$ (still in the right range to produce enough X particles from SN_{Ia}) and $T = 1\,\text{MeV}$ we find that n is already a few percents of its original density which, in this case, is a small perturbation to the cosmological energy density.

B. Case (B)

For $M \approx 100\,\text{TeV}$ and $m_{\eta_h,\omega_h} \sim 1.5, 2\,\text{MeV}$, we get $\tau_{\eta_h} \sim 10^{31}\,\text{s}$ and $\tau_{\omega_h} \sim 10^{11}\,\text{s}$. Thus, given our preceding discussion, we see that ω_h can easily be a good candidate for particle X, responsible for the SN_{cc}-generated 511-keV signal. Note also that τ_{η_h} is consistent with

[2] A low temperature mechanism for baryogenesis [34, 35] could provide the baryon asymmetry of the universe.

511-keV flux bounds [33], even if η_h is as abundant as DM. However, in this case η_h will be very long-lived and could upset the standard picture of cosmology if its relic density is too large. In case (B), the hidden sector decouples from the SM in the early universe at a temperature $T_d \approx 13\,\text{GeV}$, where we have assumed $g_* \simeq 86$. Again, the simplest assumption is that inflation resulted in a low re-heat temperature of order T_d (within the SM sector) that did not lead to any significant production of long-lived hidden hadrons. Nonetheless, given the longevity of η_h, it is interesting to see whether there is a reasonable scenario in which this particle is the dominant DM of the universe. We will examine this possibility next.

Let us assume that the primordial SM and hidden sectors evolve in a decoupled fashion, as discussed above, such that they do not come into thermal equilibrium. In this case, once the hidden sector's temperature T_h falls below $T_* \sim 1\,\text{MeV}$, its constituents hadronize. All the heavier hidden hadrons will quickly annihilate down to η_h which will then redshift as matter. Let us denote the temperature of the SM sector at this point by T_i. In order for matter-radiation equality to take place near its standard temperature of $T_f \sim 3\,\text{eV}$, we then demand

$$T_*^4 (T_f/T_i)^3 \sim T_f^4\,, \tag{17}$$

which yields

$$10\,\text{MeV} \lesssim T_i \lesssim 10^2\,\text{MeV} \quad \text{for} \quad 0.3\,\text{MeV} \lesssim T_* \lesssim 1\,\text{MeV}. \tag{18}$$

We note that the hidden dynamical scale corresponding to T_* can be somewhat lower than the hadron masses (as in QCD). Also, over the above range of T_i, the SM and hidden sectors stay decoupled, as $T_d \approx 13\,\text{GeV}$. Hence, in this setup η_h can be a realistic DM candidate, as long as the conditions (18) are met. To complete this discussion, we then suggest a scenario in which these conditions can be realized. Imagine that the universe went through a period of high scale inflation with a reheat temperature $T_r \gg T_d$, giving rise to a thermalized plasma of SM and hidden species. However, a late and milder inflation can cool this plasma down to $T \sim 1\,\text{MeV}$. As long as the late inflaton only decays into SM degrees of freedom and leads to a reheated SM plasma at $T = T_i$, consistent with (18), η_h will survive as a viable DM particle.

VI. DISCUSSION AND SUMMARY

We examined the conditions under which the 511 keV signal could be related to supernova production of an MeV-scale composite state X of hidden dynamics in simple models. Our proposed setups naturally avoid the escape fraction problem of conventional $\mathrm{SN_{Ia}}$ scenarios. However, the models we considered face some challenges, from cosmology or in relation to the spatial distribution of the signal. Nonetheless, these proposals offer interesting examples of how hidden valley dynamics could affect astrophysical observations, such as that of the 511 keV signal. Our analysis is done at a semi-quantitative level and a much more detailed study is required in order to make definite numerical statements.

Nevertheless, our study demonstrates, in a concrete manner, that hidden valley dynamics could play an important role in cosmology and astrophysics, while also having implications for discoveries at accelerators. We considered how $\mathrm{SN_{Ia}}$ [case (A)] and $\mathrm{SN_{cc}}$ [case (B)] explosions can generate new composite states with relevance to the 511 keV data.

Case (A): The $\mathrm{SN_{Ia}}$ explosions tend to be concentrated mostly near the central part of the Galaxy, where the signal originates. Conventional explanations based on $\mathrm{SN_{Ia}}$ suffer from the small escape fraction of e^+ generated in radioactive decays of the ejecta. However, in our scenario, the X particles have a large enough lifetime to release the e^+e^- decay products away from the relatively slow moving explosion front. Boltzmann suppressed production of X ensures a non-relativistic emission process and leads to a soft e^+ flux. A minimal setup requires a state Φ mediating between the hidden and SM sectors, with a mass $m_\Phi \sim$ 300 GeV, weak coupling to e^+e^- and strong coupling to hidden f-quarks. In this case, Φ may be directly discovered at a high luminosity e^+e^- TeV-scale collider. In addition, in this case, it is interesting that both type Ia and core collapse supernova dynamics are modified which may lead to other future signals in precision measurements. If the above effect has some z dependence it might lead to systematic effects for the next generation supernova observatories. Economical hidden valley models can capture the generic requirements to produce the 511 keV signal and lead to acceptable cosmic scenarios, though a more detailed analysis may be required to examine the latter.

Case (B): An alternative possibility is offered by $\mathrm{SN_{cc}}$. These explosions are characterized by temperatures far above m_X and can easily provide the required X flux. The X production is through hidden sector quark jet fragmentation. This results in a soft spectrum for the

final state e^+e^-, as required. The requisite hidden sector production is governed by a large suppression scale $m_\Phi \sim 100$ TeV, well beyond foreseeable collider reaches, which leads to long-lived states. In this case, simple hidden valley models may potentially include a DM candidate. Given that most $\mathrm{SN_{cc}}$ explosions take place in the Galactic disk outside the bulge, the e^+e^- produced in this case must be transported away from the source. Achieving the requisite transfer of e^+ away from the disk over such length scales may be a challenge within a conventional setup. However, in our scenario, the implied decay length of X is typically several kpc, allowing for an efficient deposition of e^+e^- far from the disk. This leads to a reduced disk surface brightness by diffusing the e^+e^- over a larger volume and may help transport of e^+ into the bulge, via large scale magnetic fields of the Galaxy. These effects could help explain the B/D ratio inferred from observations. However, the relevance of case (B) to the 511 keV signal depends on whether the aforementioned transport mechanism is operative in the Galaxy.

Acknowledgments

We thank M. Strassler for many helpful early discussions, collaboration on topics that led to this project, and comments on a draft version of the present paper. We also thank K. Blum, A. Gal-Yam, U. Karshon, C. Lunardini and T. McElmurry for conversations and comments. The work of H.D. is supported by the United States Department of Energy under Grant Contract DE-AC02-98CH10886. G.P. is supported by the Israel Science Foundation (grant #1087/09), EU-FP7 Marie Curie, IRG fellowship and the Peter & Patricia Gruber Award.

---

[1] J. Knodlseder *et al.*, Astron. Astrophys. **441**, 513 (2005) [arXiv:astro-ph/0506026].

[2] For a summary of prior measuremensts, see: B. J. Teegarden *et al.*, Astrophys. J. **621**, 296 (2005) [arXiv:astro-ph/0410354].

[3] E. Kalemci, S. E. Boggs, P. A. Milne and S. P. Reynolds, Astrophys. J. **640**, L55 (2006) [arXiv:astro-ph/0602233].

[4] N. Prantzos, Astron. Astrophys. **449**, 869 (2006) [arXiv:astro-ph/0511190].

[5] C. Boehm, D. Hooper, J. Silk, M. Casse and J. Paul, Phys. Rev. Lett. **92**, 101301 (2004) [arXiv:astro-ph/0309686]; D. Hooper and K. M. Zurek, Phys. Rev. D **77**, 087302 (2008) [arXiv:0801.3686 [hep-ph]].

[6] D. P. Finkbeiner and N. Weiner, Phys. Rev. D **76**, 083519 (2007) [arXiv:astro-ph/0702587].

[7] N. Arkani-Hamed, D. P. Finkbeiner, T. R. Slatyer and N. Weiner, Phys. Rev. D **79**, 015014 (2009) [arXiv:0810.0713 [hep-ph]].

[8] A. Dar, J. Goodman and S. Nussinov, Phys. Rev. Lett. **58**, 2146 (1987) [Erratum-ibid. **59**, 1871 (1987)].

[9] J. C. Higdon, R. E. Lingenfelter and R. E. Rothschild, Astrophys. J. **698**, 350 (2009) [arXiv:0711.3008 [astro-ph]]; R. E. Lingenfelter, J. C. Higdon and R. E. Rothschild, Phys. Rev. Lett. **103**, 031301 (2009) [arXiv:0904.1025 [astro-ph.HE]].

[10] P. Jean *et al.*, Astron. Astrophys. **445**, 579 (2006) [arXiv:astro-ph/0509298].

[11] M. J. Strassler and K. M. Zurek, Phys. Lett. B **651**, 374 (2007) [arXiv:hep-ph/0604261].

[12] B. Holdom, Phys. Lett. B **166**, 196 (1986). For recent applications of this mechanism to theories of dark matter also see M. Pospelov, A. Ritz and M. B. Voloshin, Phys. Lett. B **662**, 53 (2008) [arXiv:0711.4866 [hep-ph]] and Ref. [7].

[13] E. Witten, Nucl. Phys. B **149**, 285 (1979).

[14] K. Nomoto, F. K. Thielemann and K. Yokoi, Astrophys. J. **286**, 644 (1984).

[15] *Stars as Laboratories for Fundamental Physics*, Georg G. Raffelt, The University of Chicago Press, 1996.

[16] J. A. Grifols, E. Masso and S. Peris, Mod. Phys. Lett. A **4**, 311 (1989).

[17] [LEP Collaborations: ALEPH, DELPHI, L3, OPAL, LEP Electroweak Working Group, the SLD Electroweak and Heavy Flavor Groups], arXiv:hep-ex/0312023.

[18] A. Heister *et al.* [ALEPH Collaboration], Eur. Phys. J. C **28**, 1 (2003).

[19] J. Abdallah *et al.* [DELPHI Collaboration], Eur. Phys. J. C **38**, 395 (2005) [arXiv:hep-ex/0406019].

[20] M. Jaffre [CDF Collaboration and D0 Collaboration], arXiv:0909.2979 [hep-ex].

[21] H. Davoudiasl, R. Kitano, G. D. Kribs and H. Murayama, Phys. Rev. D **71**, 113004 (2005) [arXiv:hep-ph/0502176].

[22] For a discussion of the theoretical and experimental status of $(g-2)_\mu$ phenomenology see, for example, M. Passera, W. J. Marciano and A. Sirlin, Phys. Rev. D **78**, 013009 (2008)

[arXiv:0804.1142 [hep-ph]]; arXiv:1001.4528 [hep-ph].

[23] See, for example, *An Introduction to Quantum Field Theory*, M. E. Peskin and D. V. Schroeder, Addison-Wesley Publishing Company, 1995.

[24] N. Arkani-Hamed, S. Dimopoulos and G. R. Dvali, Phys. Lett. B **429**, 263 (1998) [arXiv:hep-ph/9803315].

[25] G. F. Giudice, R. Rattazzi and J. D. Wells, Nucl. Phys. B **544**, 3 (1999) [arXiv:hep-ph/9811291].

[26] E. A. Mirabelli, M. Perelstein and M. E. Peskin, Phys. Rev. Lett. **82**, 2236 (1999) [arXiv:hep-ph/9811337].

[27] E. L. Chupp, W. T. Vestrand and C. Reppin, Phys. Rev. Lett. **62**, 505 (1989).

[28] C. Amsler *et al.* [Particle Data Group], Phys. Lett. B **667**, 1 (2008).

[29] W. Braunschweig *et al.* [TASSO Collaboration], Z. Phys. C **47**, 187 (1990).

[30] We thank Avishay Gal-Yam for comments on this point.

[31] A. Mirizzi, G. G. Raffelt and P. D. Serpico, JCAP **0605**, 012 (2006) [arXiv:astro-ph/0604300].

[32] R. Lazauskas, C. Lunardini and C. Volpe, JCAP **0904**, 029 (2009) [arXiv:0901.0581 [astro-ph.SR]].

[33] C. Picciotto and M. Pospelov, Phys. Lett. B **605**, 15 (2005) [arXiv:hep-ph/0402178].

[34] I. Affleck and M. Dine, Nucl. Phys. B **249**, 361 (1985).

[35] S. Dimopoulos and L. J. Hall, Phys. Lett. B **196**, 135 (1987).

[36] S. Hannestad, Phys. Rev. D **70**, 043506 (2004) [arXiv:astro-ph/0403291]; G. B. Gelmini, arXiv:0805.1824 [hep-ph].